QCD factorisation and flavour symmetries illustrated in $B_{d,s} \to KK$ decays

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We present a new analysis of $B_{d,s} \to KK$ modes within the SM, relating them in a controlled way through $SU(3)$-flavour symmetry and QCD-improved factorisation. We propose a set of sum rules for $B_{d,s} \to K^0\bar{K}^0$ observables. We determine $B_s \to KK$ branching ratios and CP-asymmetries as functions of $A_{dir}(B_d \to K^0\bar{K}^0)$, with a good agreement with current experimental measurements of CDF. Finally, we predict the amount of U-spin breaking between $B_d \to \pi^+\pi^-$ and $B_s \to K^+K^-$.

The current data in B-physics suggests that B_d-decays agree well with SM predictions, while B_s-decays remain poorly known and might be affected by New Physics. Within the Standard Model, the CKM mechanism correlates the electroweak part of these transitions, but quantitative predictions are difficult due to hadronic effects. The latter can be estimated relying on the approximate $SU(3)$-flavour symmetry of QCD : information on hadronic effects, extracted from data in one channel, can be exploited in other channels related by flavour symmetry, leading to more accurate predictions within the Standard Model.

In addition to isospin symmetry, an interesting theoretical tool is provided by U-spin symmetry, which relates d- and s-quarks. Indeed, this symmetry holds for long- and short-distances and does not suffer from electroweak corrections, making it a valuable instrument to analyse processes with significant penguins and thus a potential sensitivity to New Physics. However, due to the significant difference $m_s - m_d$, U-spin breaking corrections of order 30 % may occur, depending on the processes.

As a first application of U-spin, relations were obtained between $B_d \to \pi^+\pi^-$ and $B_s \to K^+K^-$. This led to correlations among the observables in the two decays such as branching ratios and CP asymmetries [1, 2] and to a prediction for $Br(B_s \to K^+K^-) = (35^{+73}_{-20}) \cdot 10^{-6}$ [3]. These results helped to investigate the potential of such decays to discover New Physics [4, 5]. Unfortunately, the accuracy of the method is limited not only by the persistent discrepancy between Babar and Belle on $B_d \to \pi^+\pi^-$ CP asymmetries, but also by poorly known U-spin corrections. In these analyses, the ratio of tree contributions $R_c = |T^s_{K\pm}/T^d_{\pi\pm}|$ was taken from QCD sum rules as 1.76 ± 0.17, updated to $1.52^{+0.18}_{-0.14}$ [6]. In addition, the ratio of penguin-to-tree ratios $\xi = |(P^s_{K\pm}/T^s_{K\pm})/(P^d_{\pi\pm}/T^d_{\pi\pm})|$ was assumed equal to 1 [3] or 1 ± 0.2 [4, 5] in agreement with rough estimates within QCD factorisation (QCDF) [7]. Recent updates on U-spin methods were given during this workshop [8].

QCDF may complement flavour symmetries by a more accurate study of short-distance effects. However, this expansion in α_s and $1/m_b$ cannot predict some significant $1/m_b$-suppressed long-distance effects, which have to be estimated through models. Recently, it was proposed to combine QCDF and U-spin in the decays mediated by penguin operators $B_d \to K^0\bar{K}^0$ and $B_s \to K^0\bar{K}^0$ [9].

The SM amplitude for a B decaying into two mesons can be split into tree and penguin contributions [10]:

$$\bar{A} \equiv A(\bar{B}_q \to M\bar{M}) = \lambda_u^{(q)} T_M^{qC} + \lambda_c^{(q)} P_M^{qC} , \qquad (1)$$

with C denoting the charge of the decay products, and the products of CKM factors $\lambda_p^{(q)} = V_{pb}V^*_{pq}$. Using QCDF [11, 12], one can perform a $1/m_b$-expansion of the amplitude. The tree and penguin contributions in $\bar{B}_s \to K^+K^-$ and $\bar{B}_s \to K^0\bar{K}^0$ in QCDF are :

$$\begin{aligned}
\hat{T}^{s\,\pm} &= \bar\alpha_1 + \bar\beta_1 && (2)\\
&\quad + \bar\alpha_4^u + \bar\alpha^u_{4EW} + \bar\beta^u_3 + 2\bar\beta^u_4 - \frac{1}{2}\bar\beta^u_{3EW} + \frac{1}{2}\bar\beta^u_{4EW} \\
\hat{P}^{s\,\pm} &= \bar\alpha_4^c + \bar\alpha^c_{4EW} + \bar\beta^c_3 + 2\bar\beta^c_4 - \frac{1}{2}\bar\beta^c_{3EW} + \frac{1}{2}\bar\beta^c_{4EW} && (3)\\
\hat{T}^{s\,0} &= \bar\alpha_4^u - \frac{1}{2}\bar\alpha^u_{4EW} + \bar\beta^u_3 + 2\bar\beta^u_4 - \frac{1}{2}\bar\beta^u_{3EW} - \bar\beta^u_{4EW} && (4)\\
\hat{P}^{s\,0} &= \bar\alpha_4^c - \frac{1}{2}\bar\alpha^c_{4EW} + \bar\beta^c_3 + 2\bar\beta^c_4 - \frac{1}{2}\bar\beta^c_{3EW} - \bar\beta^c_{4EW} && (5)
\end{aligned}$$

where $\hat{P}^{sC} = P^{sC}/A^s_{KK}$, $\hat{T}^{sC} = T^{sC}/A^s_{KK}$ and $A^q_{KK} = M^2_{B_q} F_0^{\bar{B}_q \to K}(0) f_K G_F/\sqrt{2}$. The superscripts identify the channel and the bar denotes quantities for decays with a spectator s-quark. The tree and penguin contributions $T^{d\,0}$ and $P^{d\,0}$ for $\bar{B}_d \to K^0\bar{K}^0$ have the same structure as eqs. (4) and (5), with unbarred α's and β's recalling the different nature of the spectator d-quark.

At NLO in α_s, α's are linear combinations of vertex corrections, hard-spectator terms and penguin contractions, whereas β's are sums of annihilation contributions. The weights of the various contributions are expressed in terms of α_s and Wilson coefficients [12]. α's and β's contain the two most significant terms in the $1/m_b$ expansion: the LO terms, dominated by short distances, and the NLO terms in $1/m_b$ that include the potentially large long-distance corrections. The latter, parameterised in QCDF through quantities denoted X_H (in power corrections to the hard-scattering part of α_i) and X_A (in the annihilation parameters β_i), are singled out since they may upset the quick convergence of the $1/m_b$ expansion. The other $1/m_b$-suppressed contributions, dominated by short distances, are under control and small, i.e, leading to a $\mathcal{O}(5-10\%)$ error. In [9], we showed that comparing

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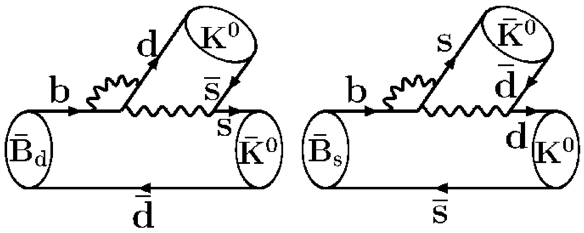

FIG. 1: Diagrams contributing to $\bar{B}_d \to K^0\bar{K}^0$ (left) and $\bar{B}_s \to K^0\bar{K}^0$ (right) related through U-spin transformations.

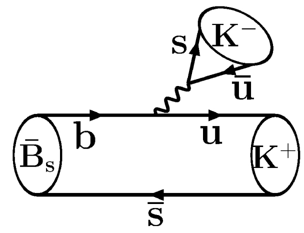

FIG. 2: Diagram contributing to $\bar{B}_s \to K^+\bar{K}^-$ from a tree operator, without counterpart in $\bar{B}_d \to K^0\bar{K}^0$.

B_d- and B_s-decays into the same final states helps to cancel the potentially large long-distance $1/m_b$-suppressed effects ($X_{A,H}$), yielding improved SM predictions.

I. SUM RULES

Let us start with the difference $\Delta_d \equiv T^{d0} - P^{d0}$ is free from the troublesome NLO infrared-divergence (modelled by $X_{A,H}$) that may be enhanced numerically by the chiral factor $r_\chi^K = 2m_K^2/m_b/m_s$ from twist-3 distribution amplitudes. Hard-scattering (X_H) and annihilation (X_A) terms occur in both penguin and tree contributions, but remarkably they cancel in the short-distance difference:

$$\begin{aligned}\Delta_d &= A^d_{KK}[\alpha_4^u - \alpha_4^c + \beta_3^u - \beta_3^c + 2\beta_4^u - 2\beta_4^c] \qquad (6)\\ &= A_{KK} \times \alpha_s C_F C_1 \times [\bar{G}(m_c^2/m_b^2) - \bar{G}(0)]/(4\pi N_c)\end{aligned}$$

neglecting (small) electroweak contributions. The function $\bar{G} = G_K - r_\chi^K \hat{G}_K$ combines one-loop integrals from the penguin terms P_4 and P_6 defined in Sec 2.4 in ref. [12]. The same cancellation of long-distance $1/m_b$-corrections happens for $\Delta_s \equiv T^{s0} - P^{s0}$. Taking into account the uncertainties coming from the QCDF inputs [12], we get $\Delta_d = (1.09 \pm 0.43) \cdot 10^{-7} + i(-3.02 \pm 0.97) \cdot 10^{-7}$GeV and $\Delta_s = (1.03 \pm 0.41) \cdot 10^{-7} + i(-2.85 \pm 0.93) \cdot 10^{-7}$GeV.

These two theoretical quantities can be related to observables, namely the corresponding branching ratio and coefficients of the time-dependent CP-asymmetry:

$$\begin{aligned}&\frac{\Gamma(B_d(t) \to K^0\bar{K}^0) - \Gamma(\bar{B}_d(t) \to K^0\bar{K}^0)}{\Gamma(B_d(t) \to K^0\bar{K}^0) + \Gamma(\bar{B}_d(t) \to K^0\bar{K}^0)} \qquad (7)\\ &\qquad = \frac{A^{d0}_{dir}\cos(\Delta M \cdot t) + A^{d0}_{mix}\sin(\Delta M \cdot t)}{\cosh(\Delta\Gamma_d t/2) - A^{d0}_{\Delta}\sinh(\Delta\Gamma_d t/2)},\end{aligned}$$

where we define [1]: $A^{d0}_{dir} = (|A|^2 - |\bar{A}|^2)/(|A|^2 + |\bar{A}|^2)$, $A^{d0}_{\Delta} + iA^{d0}_{mix} = -(2e^{-i\phi_d}A^*\bar{A})/(|A|^2 + |\bar{A}|^2)$ and ϕ_d the phase of $B_d - \bar{B}_d$ mixing. A^{d0}_{Δ} is unlikely to be measured due to the small width difference $\Delta\Gamma_d$, but it can be obtained from the other asymmetries by means of the relation $|A^{d0}_{\Delta}|^2 + |A^{d0}_{dir}|^2 + |A^{d0}_{mix}|^2 = 1$.

One can derive the following relation for $B_d \to K^0\bar{K}^0$:

$$\begin{aligned}|\Delta_d|^2 &= \frac{BR^{d0}}{L_d}\{x_1 + [x_2 \sin\phi_d - x_3\cos\phi_d]A^{d0}_{mix} \qquad (8)\\ &\qquad - [x_2\cos\phi_d + x_3\sin\phi_d]A^{d0}_{\Delta}\},\end{aligned}$$

where $L_d = \tau_d\sqrt{M_{Bd}^2 - 4M_K^2}/(32\pi M_{Bd}^2)$ and:

$$\begin{aligned}x_1 &= [|\lambda_c^{(d)}|^2 + |\lambda_u^{(d)}|^2 - 2|\lambda_c^{(d)}||\lambda_u^{(d)}|\cos\gamma]/n^2,\\ x_2 &= -[|\lambda_c^{(d)}|^2 + |\lambda_u^{(d)}|^2\cos 2\gamma - 2|\lambda_c^{(d)}||\lambda_u^{(d)}|\cos\gamma]/n^2,\\ x_3 &= -[1 - \cos\gamma \times |\lambda_u^{(d)}|/|\lambda_c^{(d)}|]/n,\end{aligned}$$

with $n = 2|\lambda_c^{(d)}||\lambda_u^{(d)}|\sin\gamma$. A similar relation between Δ_s and $B_s \to K^0\bar{K}^0$ observables is obtained by replacing $|\lambda_u^{(d)}| \to |\lambda_u^{(s)}|$, $|\lambda_c^{(d)}| \to -|\lambda_c^{(s)}|$, and $d \to s$ for all indices.

These sum rules can be used either as a way to extract the SM value of one observable (say $|A^{s0}_{dir}|$) in terms of the two others (BR^{s0} and A^{s0}_{mix}) and Δ_s, as a SM consistency test between BR^{s0}, $|A^{s0}_{dir}|$ and A^{s0}_{mix} (and similarly for the $B_d^0 \to K^0\bar{K}^0$ observables), or as a way of determining CKM parameters [13]. These relations are free from the long-distance power-suppressed model-dependent quantities X_A and X_H that are a main error source in the direct computation of A^{s0}_{dir} within QCDF.

II. FLAVOUR SYMMETRIES AND QCDF

Using U-spin symmetry, we can relate the two penguin-mediated decays $\bar{B}_d \to K_0\bar{K}_0$ and $\bar{B}_s \to K_0\bar{K}_0$, as exemplified in fig. 1 (see also ref. [14] in relation to $B \to \pi\pi$). U-spin breaking should be much smaller here than usual: it does not affect final-state interaction since both decays involve the same outgoing state, and it shows up mainly in power-suppressed effects. This is confirmed by QCDF:

$$\begin{aligned}P^{s0} &= fP^{d0}\Big[1 + (A^d_{KK}/P^{d0})\Big\{\delta\alpha_4^c - \delta\alpha^c_{4EW}/2\\ &\qquad + \delta\beta_3^c + 2\delta\beta_4^c - \delta\beta^c_{3EW}/2 - \delta\beta^c_{4EW}\Big\}\Big],\\ T^{s0} &= fT^{d0}\Big[1 + (A^d_{KK}/T^{d0})\Big\{\delta\alpha_4^u - \delta\alpha^u_{4EW}/2 \qquad (9)\\ &\qquad + \delta\beta_3^u + 2\delta\beta_4^u - \delta\beta^u_{3EW}/2 - \delta\beta^u_{4EW}\Big\}\Big],\end{aligned}$$

where we define the U-spin breaking differences $\delta\alpha_i^p \equiv \bar{\alpha}_i^p - \alpha_i^p$ (id. for β). Apart from the factorisable ratio :

$$f = A^s_{KK}/A^d_{KK} = M^2_{B_s}F_0^{\bar{B}_s\to K}(0)/[M^2_{B_d}F_0^{\bar{B}_d\to K}(0)]$$

which should be computed on the lattice, U-spin breaking arises through $1/m_b$-suppressed contributions in which most long-distance contributions have cancelled out.

First, the hard-spectator scattering ($\delta\alpha$) probes the difference between B_d- and B_s-distribution amplitudes which is expected small, since the dynamics of the heavy-light meson in the limit $m_b \to \infty$ should vary little from B_d and B_s. Second, the annihilation contributions ($\delta\beta$) contain a U-spin breaking part when the gluon is emitted from the light quark in the $B_{d,s}$-meson (this effect from A_1^i and A_2^i defined in [12] is neglected in the QCDF model for annihilation terms). Taking the hadronic parameters in [12], we obtain $|P^{s0}/(fP^{d0}) - 1| \leq 3\%$ and $|T^{s0}/(fT^{d0}) - 1| \leq 3\%$.

Relations exist between $\bar{B}_d \to K_0\bar{K}_0$ and $\bar{B}_s \to K^+K^-$ as well. A combination of U-spin and isospin rotations leads from the penguin contribution in $\bar{B}_d \to K_0\bar{K}_0$ to that in $\bar{B}_s \to K_0\bar{K}_0$, then to $\bar{B}_s \to K^+K^-$, up to electroweak corrections (it corresponds to fig. 1 up to replacing $d \to u$ in the right-hand diagram). On the other hand, there are no such relations between trees, since $\bar{B}_s \to K^+K^-$ contains tree contributions (see fig. 2) which have no counterpart in the penguin-mediated decay $\bar{B}_d \to K_0\bar{K}_0$. This is seen in QCDF as well:

$$P^{s\pm} = fP^{d0}\Big[1 + \frac{A^d_{KK}}{P^{d0}}\Big\{\frac{3}{2}(\alpha^c_{4EW} + \beta^c_{4EW}) + \delta\alpha^c_4 + \delta\alpha^c_{4EW} + \delta\beta^c_3 + 2\delta\beta^c_4 - \frac{1}{2}(\delta\beta^c_{3EW} - \delta\beta^c_{4EW})\Big\}\Big], \quad (10)$$

$$\frac{T^{s\pm}}{A^s_{KK}\bar{\alpha}_1} = 1 + \frac{T^{d0}}{A^d_{KK}\bar{\alpha}_1} + \frac{1}{\bar{\alpha}_1}\Big\{\bar{\beta}_1 + \frac{3}{2}(\alpha^u_{4EW} + \beta^u_{4EW}) + \delta\alpha^u_4 + \delta\alpha^u_{4EW} + \delta\beta^u_3 + 2\delta\beta^u_4 - \frac{1}{2}(\delta\beta^u_{3EW} - \delta\beta^u_{4EW})\Big\}.$$

Terms are ordered in decreasing size (in particular, curly brackets in $T^{s\pm}$ should be tiny). From QCDF, we obtain the following bounds: $|P^{s\pm}/(fP^{d0}) - 1| \leq 2\%$ and $|T^{s\pm}/(A^s_{KK}\bar{\alpha}_1) - 1 - T^{d0}/(A^d_{KK}\bar{\alpha}_1)| \leq 4\%$. The latter shows that flavour-symmetry breaking corrections are smaller than $T^{d0}/(A^d_{KK}\bar{\alpha}_1) = O(10\%)$. Fortunately, $T^{s\pm}$ is strongly CKM suppressed in $B_s \to K^+K^-$ so that the uncertainty on its QCDF determination will affect the branching ratio and CP-asymmetries only marginally.

III. SM PREDICTIONS FOR $B_s \to KK$ DECAYS

The dynamics of $B_d \to K^0\bar{K}^0$ involves three hadronic real parameters (modulus of the tree, modulus of the penguin and relative phase) which we can pin down through three observables: BR^{d0}, A^{d0}_{dir} and A^{d0}_{mix}. Only $BR^{d0} = (0.96 \pm 0.25) \cdot 10^{-6}$ [16] has been measured. However the direct asymmetry A^{d0}_{dir} should be observable fairly easily (for instance, $A^{d0}_{dir} = 0.19 \pm 0.06$ in QCDF) whereas the mixed asymmetry is likely small ($A^{d0}_{mix} = 0.05 \pm 0.05$ in QCDF).

If only A^{d0}_{dir} becomes available, we have only 2 experimental constraints for 3 hadronic parameters. Then we may exploit a theoretically well-controlled QCDF constraint to get T^{d0} and P^{d0} from BR^{d0}, A^{d0}_{dir} and the QCDF value of $\Delta_d \equiv T^{d0} - P^{d0}$, free from infrared divergences and thus with little model dependence. This system yields two constraints in the complex plane (x_P, y_P) for P^{d0} : a circular ring and a diagonal strip [9], which can be satisfied only if $|A^{d0}_{dir}| < 0.2$, and then yield two different solutions with opposite signs for Im P^{d0}, yielding two solutions for (P^{d0}, T^{d0}).

From the measured value of the branching ratio for $B_d \to K^0\bar{K}^0$, and choosing a particular value of the direct asymmetry A^{d0}_{dir}, we get the penguin and tree contributions as explained above. Then, the bounds in II yield the hadronic parameters in $B_s \to KK$ decays up to small uncertainties. To be more conservative, we actually stretch the bounds in II relating B_d and B_s hadronic parameters up to 5 % in order to account for well-behaved short-distance $1/m_b$-corrections not yet included.

We obtain observables as functions of A^{d0}_{dir} in Table I. In the case of the branching ratios, we have split the error in two parts. The first uncertainty comes from the QCDF estimates of Δ_d and $\bar{\alpha}_1$, the theoretical constraints derived in II to relate B_d and B_s decays and the measurement of BR^{d0} (this experimental uncertainty dominates the others). The second error stems from (factorisable) U-spin breaking terms: $f = 0.94 \pm 0.2$ (cf. [12]).

Table I corresponds only to the solution of the constraints with Im $P^{d0} > 0$. But BR^{d0}, A^{d0}_{dir} and Δ_d yield two different solutions for (T^{d0}, P^{d0}), and thus for $(T^{s\pm}, P^{s\pm})$. Only one solution is physical, whereas the other stems from the non-linear nature of the constraints. We can lift this ambiguity by exploiting a channel related to $B_s \to K^+K^-$ through U-spin, namely $\bar{B}_d \to \pi^+\pi^-$ [1, 3, 4, 5, 17]. As explained in ref. [9], this allows us to reject one of the two solutions, and to compute for the other the U-spin breaking parameters :

$$R_C = \left|\frac{T^{s\pm}}{T^{d\pm}_{\pi\pi}}\right| = 2.0 \pm 0.6 \qquad \xi = \left|\frac{P^{s\pm}}{T^{s\pm}}\frac{P^{d\pm}_{\pi\pi}}{T^{d\pm}_{\pi\pi}}\right| = 0.8 \pm 0.3 \quad (11)$$

The determination of $BR^{s\pm}$ is improved compared to the U-spin extraction from $\bar{B}_d \to \pi^+\pi^-$ [3, 4, 5], and in good agreement with the recent CDF measurement [18]:

$$BR^{s\pm}\big|_{exp} \cdot 10^6 \;=\; 24.4 \pm 1.4 \pm 4.6 \quad (12)$$

IV. CONCLUSIONS

We have combined experimental data, flavour symmetries and QCDF to propose sum rules for $B_{d,s} \to K^0\bar{K}^0$ observables and to give SM constraints on $B_s \to K\bar{K}$ in a controlled way. Tree (T^{d0}) and penguin (P^{d0}) contributions to $B_d \to K^0\bar{K}^0$ can be determined by combining the currently available data with $|T^{d0} - P^{d0}|$, which can be accurately computed in QCDF because long-distance effects, seen as infrared divergences, cancel in this difference. U-spin suggests accurate relations between these hadronic parameters in $B_d \to K^0\bar{K}^0$ and those in $B_s \to K^0\bar{K}^0$. Actually, we expect similar long-distance effects since the $K^0\bar{K}^0$ final state is invariant

	$BR^{s0} \times 10^6$	$A^{s0}_{dir} \times 10^2$	$A^{s0}_{mix} \times 10^2$	$BR^{s\pm} \times 10^6$	$A^{s\pm}_{dir} \times 10^2$	$A^{s\pm}_{mix} \times 10^2$
$A^{d0}_{dir} = -0.2$	$18.4 \pm 6.5 \pm 3.6$	0.8 ± 0.3	-0.3 ± 0.8	$21.9 \pm 7.9 \pm 4.3$	24.3 ± 18.4	24.7 ± 15.5
$A^{d0}_{dir} = -0.1$	$18.2 \pm 6.4 \pm 3.6$	0.4 ± 0.3	-0.7 ± 0.7	$19.6 \pm 7.3 \pm 4.2$	35.7 ± 14.4	7.7 ± 15.7
$A^{d0}_{dir} = 0$	$18.1 \pm 6.3 \pm 3.6$	0 ± 0.3	-0.8 ± 0.7	$17.8 \pm 6.0 \pm 3.7$	37.0 ± 12.3	-9.3 ± 10.6
$A^{d0}_{dir} = 0.1$	$18.2 \pm 6.4 \pm 3.6$	-0.4 ± 0.3	-0.7 ± 0.7	$16.4 \pm 5.7 \pm 3.3$	29.7 ± 19.9	-26.3 ± 15.6
$A^{d0}_{dir} = 0.2$	$18.4 \pm 6.5 \pm 3.6$	-0.8 ± 0.3	-0.3 ± 0.8	$15.4 \pm 5.6 \pm 3.1$	6.8 ± 28.9	-40.2 ± 14.6

TABLE I: Observables for $\bar{B}_s \to K^0 \bar{K}^0$ and $\bar{B}_s \to K^+K^-$ as functions of the direct asymmetry $A_{dir}(\bar{B}_d \to K^0\bar{K}^0)$ within the SM. We take $\lambda_u^{(d)} = 0.0038 \cdot e^{-i\gamma}$, $\lambda_c^{(d)} = -0.0094$, $\lambda_u^{(s)} = 0.00088 \cdot e^{-i\gamma}$, $\lambda_c^{(s)} = 0.04$, and $\gamma = 62°$, $\phi_d = 47°$, $\phi_s = -2°$ [15].

under the d-s exchange. Short distances are also related since the two processes are mediated by penguin operators through diagrams with the same topologies. U-spin breaking arises only in a few places : factorisable corrections encoded in $f = [M^2_{B_s} F^{\bar{B}_s \to K}(0)]/[M^2_{B_d} F^{B_d \to K}(0)]$, and non-factorisable corrections from weak annihilation and spectator scattering.

Because of these expected tight relations, QCDF can be relied upon to assess U-spin breaking between the two decays. Indeed, up to the factorisable factor f, penguin (as well as tree) contributions to both decays are numerically very close. Penguins in $B_d \to K^0\bar{K}^0$ and $B_s \to K^+K^-$ should have very close values as well, whereas no such relation exists for the (CKM-suppressed) tree contribution to the latter.

These relations among hadronic parameters, inspired by U-spin considerations and quantified within QCD factorisation, can be exploited to predict : $Br(B_s \to K^0\bar{K}^0) = (18 \pm 7 \pm 4 \pm 2) \cdot 10^{-6}$ and $Br(B_s \to K^+K^-) = (20 \pm 8 \pm 4 \pm 2) \cdot 10^{-6}$, in very good agreement with the latest CDF measurement. The same method provides significantly improved determinations of the U-spin breaking ratios $\xi = 0.8 \pm 0.3$ and $R_c = 2.0 \pm 0.6$. These results have been exploited to determine the impact of supersymmetric models on these decays [19].

Our method merges ingredients from flavour symmetries and QCD factorisation in order to improve the accuracy of the predictions. Flavour symmetries must rely on global fudge factors typically of order 30%, without providing hints where symmetry breaking is large or small (typically, ξ is guesstimated, and not computed). QCD factorisation allows one to classify the contributions of the various operators of the effective Hamiltonian, but it suffers from a model-dependence in potentially large $1/m_b$ corrections. We use QCD factorisation to determine the short-distance contributions, and we replace models for long distances by experimental pieces of information from other decays related by flavour symmetries. Both methods are exploited optimally to yield accurate predictions for $B_s \to KK$ as functions of B_d observables.

If sizeable NP effects occur, the SM correlations between B_d and B_s decays exploited here should be broken, leading to departure from our predictions. New results on $B \to K$ form factors and on the $B_d \to K^0\bar{K}^0$ branching ratio and direct CP-asymmetry should lead to a significant improvement of the SM predictions in the B_s sector. The potential of other pairs of nonleptonic B_d and B_s decays remains to be investigated.

Acknowledgments

I thank J. Matias and J. Virto for a very enjoyable collaboration. This work was supported in part by the EU Contract No. MRTN-CT-2006-035482, "FLAVIAnet".

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